

On the thermalization of quarkonia at the LHC

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We argue that the relative yields of Υ states observed at the LHC can be understood as bottomonium states coming to early thermal equilibrium and then freezing out. The bottomonium freezeout temperature is approximately 250 MeV. We examine its systematics as a function of centrality. We remark on the interesting differences seen by the CMS and ALICE experiments in the charmonium sector.

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A thermal medium can disrupt the formation of a bound state of a heavy quark, Q and antiquark, \bar{Q} [1]. This is a multi-scale problem, involving the temperature T and the quark mass M . The quark is heavy, *i.e.*, $M/T \gg 1$. For the charmonium this ratio is about 5–10 and for the bottomonium it is around 15–30; so both these flavours can be considered to be heavy in this sense. Also, $M/\Lambda \gg 1$, where $\Lambda \simeq m_\rho/2$ is the typical scale of QCD [27]. This second comparison implies that in the $\bar{Q}Q$ bound state the quarks are slow, with velocity $v^2 \simeq 0.23$ for charm and $v^2 \simeq 0.08$ for bottom [2]. In NRQCD counting, the binding energy $B \simeq Mv^2 \simeq \Lambda$ in both cases [3]. For the temperature range of relevance, we find that $B/T \simeq \Lambda/T \simeq 1$. As a result, thermal effects can drastically modify the bound state.

Since a thermal medium can be formed in PbPb collisions, but not in pp collisions, there should be certain systematic differences between the yields in these two cases [4]. Such effects were first seen at the SPS [5], but there were important backgrounds to the signal which came from the initial state via parton density effects [6] and the final state through comover interactions [7]. Furthermore, the observed effects switched on slowly with centrality and nuclear size, and so were difficult to interpret clearly [8].

At the LHC the experimental situation has changed drastically. The the 1S, 2S and 3S states of the Υ [9] and the 1S and 2S states of the J/ψ [10, 11] have been studied, and clear differences between pp and PbPb collisions have been established. For any hadron h , we will use the notation $R_{\text{PbPb}}[h] = N_{\text{PbPb}}[h]/N_{\text{pp}}[h]$ where $N_{\text{PbPb}}[h]$ is the yield of h in PbPb collisions and $N_{\text{pp}}[h]$ in pp collisions. The new experimental results show sequential suppression [12] very clearly: $R_{\text{PbPb}}[\Upsilon(1S)]$, $R_{\text{PbPb}}[\Upsilon(2S)]$ and $R_{\text{PbPb}}[\Upsilon(3S)]$ are not equal, nor are $R_{\text{PbPb}}[J/\psi]$ and $R_{\text{PbPb}}[\psi(2S)]$.

Initial state effects cannot be the explanation, since the relevant values of the Bjorken variable are almost equal for the three bottomonium states, and a different common value for the two charmonium states, implying that parton density effects would be the same. Final state comover interactions were invoked even at the LHC in order to explain the observed suppression of J/ψ [13]. The comoving material which is thermalized gives rise to the signal. The comovers responsible for the background are unthermalized and relatively cold spectators from the initial PbPb collision. However, the data are taken at central rapidity and separated from the spectator fragments by $\Delta y \simeq 6$. So comover interactions cannot be the explanation for the observations. In this cleaner environment it would be interesting to check whether the data allow a thermal interpretation.

We begin by noting that the equilibrium density of heavy quarkonia, with mass $\mathcal{M} = 2M - B \simeq M(2 - v^2)$ and fugacity z ,

$$n \simeq (\mathcal{M}T)^{3/2} \exp\left(-\frac{\mathcal{M}}{T}\right) z, \quad (1)$$

is small. As a result, the mean distance between quarkonia in thermal equilibrium, λ , is large. In fact, we find that $T\lambda = z^{-1/3} \sqrt{T/\mathcal{M}} \exp(2\mathcal{M}/3T)$. For $z = 1$ this dimensionless number is over 100 for charmonia and over 10 million for bottomonia. So, for both the charm and bottom systems, the exponential dominates, and T is much larger than the inverse of the mean inter-quarkonium spacing. However, essentially all the $\bar{Q}Q$ pairs in question come from initial hard scatterings, so z should be computed in perturbative QCD, and the values are much larger than unity [4]. Nevertheless, the hierarchy of length scales in the fireball remains $M \gg B \simeq T \gg 1/\lambda$.

This dimensional analysis indicates that the net production rate of quarkonia in thermal equilibrium may be obtained largely by understanding the changes in the spectral density of the quarkonium states in thermal equilibrium.

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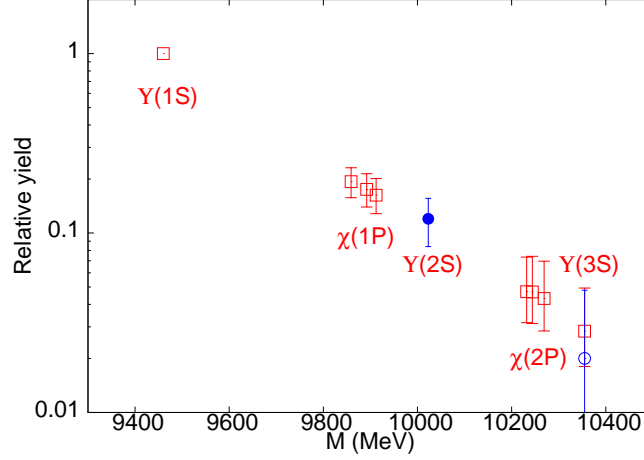


FIG. 1: Predictions for the yields of bottomonium mesons relative to the $\Upsilon(1S)$, using the simple model described in the text. The filled circle, corresponding to the relative yield of the $\Upsilon(2S)$ state, is used as an input. The rest are predictions. The observed data for the $\Upsilon(3S)$ state is shown as an unfilled circle.

Current attempts at extracting them in lattice QCD indicate that there is generally an increase in the width, Γ , of these quarkonium states at finite T [14]. This width has an interpretation as the inclusive rate of all reactions in which the quarkonium state goes into an unbound $\bar{Q}Q$ pair. In fact, understanding the phenomenon of quarkonium yields does not require a full knowledge of the spectral function; knowing the reaction rate, Γ , suffices. It was pointed out in [15] that one must add the reverse recombination processes to the screening [1] and dissociation [16] previously considered. In fact, if the heavy-quarkonium system comes to thermal equilibrium, then detailed balance requires both binding and unbinding reactions.

Defining the formation time, $\tau = 1/\Gamma$, NRQCD power counting has been used to argue that $1/\tau \simeq B = Mv^2$ [17]. This implies that $1/\tau$ is of order Λ and hence non-perturbative. However, this estimate does not involve T . Some phenomenological estimates can be found in [18, 19]. A more systematic approach is to perform a weak coupling expansion in a static limit of NRQCD at finite temperature [20, 21], from which the estimates $\Gamma \simeq g^2 T$ are obtained. However, to control weak coupling estimates, one has to have $\alpha_s \ll 1$ and hence $T \gg \Lambda$. Since we estimated that $B \simeq \Lambda \simeq T$, these weak-coupling estimates, while illuminating, may not be entirely relevant to the physical situation. Nevertheless, if one boldly extends this estimate to the region $g \simeq 1$, then $\Gamma \simeq T$, a result which is generic in effective models. This simple scaling is also consistent with recent lattice computations which indicate that for both the η_b and $\Upsilon(1S)$ one has $\Gamma = T$ (within statistical errors) immediately above T_c . [22].

If Γ for each of the resonances is large, then this would imply that the various bound and unbound levels of the $\bar{Q}Q$ system could come to thermal equilibrium with each other in a very short time. Such a situation is far simpler than that in pp collisions, where one must take into account details of how primary produced $\bar{Q}Q$ pairs hadronize into mesons. A thermalized system of this kind would stay in equilibrium until Γ falls below the expansion rate and the heavy-quark system freezes out. This resembles models of statistical hadronization [23]. However, unlike those, we do not force the heavy-quarkonia to freeze out at the same temperature as particles made with lighter quarks. Instead we let the data decide the freezeout temperature.

While the ratio R_{PbPb} is very useful for comparing pp and PbPb collisions, it is not the most appropriate quantity for examining thermal physics. Instead we examine the relative yields of different resonances. The ratios of the yields of various Υ states to that of the $\Upsilon(1S)$ are quoted in [9]—

$$\begin{aligned} Y[\Upsilon(2S)] &= \frac{N_{\text{PbPb}}[\Upsilon(2S)]}{N_{\text{PbPb}}[\Upsilon(1S)]} = 0.12 \pm 0.03 \pm 0.02 \\ Y[\Upsilon(3S)] &= \frac{N_{\text{PbPb}}[\Upsilon(3S)]}{N_{\text{PbPb}}[\Upsilon(1S)]} = 0.02 \pm 0.02 \pm 0.02 \end{aligned} \quad (2)$$

The bottomonium states were reconstructed from data taken at beam energy of 2.76 TeV per nucleon from the muons they decay into. The muon transverse momenta were greater than 4 GeV and rapidity less than 2.4. The large error bar on the second ratio means that it cannot be used as an input in the subsequent computations.

We ask whether these two pieces of data can be explained by a thermal model such as in eq. (1). Fortunately, when we consider the yield ratios Y , there is just a single parameter in this model, namely the freezeout temperature,

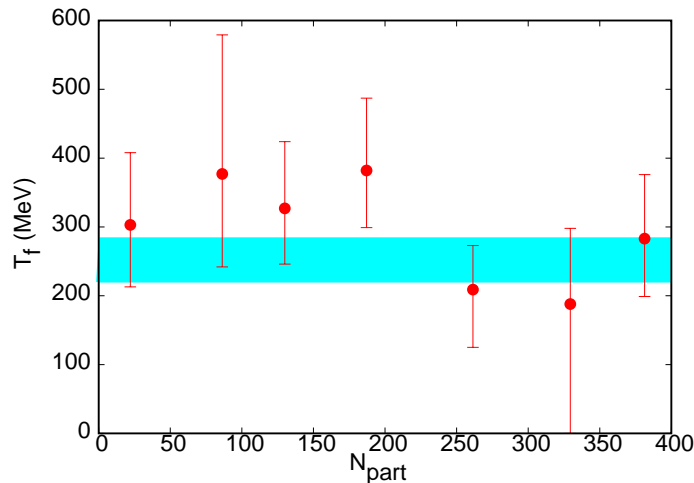


FIG. 2: The freezeout temperatures inferred from the CMS data on the centrality dependence of $Y[\Upsilon(2S)]$. The band shows the result quoted in eq. (3). Within the large error bars shown it is not possible to say with certainty whether T_f is larger for more peripheral collisions.

T_f . However, one cannot just insert the values in eq. (2) into eq. (1), and invert it to obtain the temperature. One must take into account the fact that after the mesons freeze out, the excited states may decay into the lower-lying quarkonium states before reaching the detector. This modifies the primordial densities. This has to be corrected for when extracting T_f . It turns out that inverting eq. (1) gives a reasonably close approximation, and an efficient method for extracting T_f is an iterative improvement over this guess. Adding the errors in eq. (2) in quadrature, and neglecting the errors in the branching ratios, we find

$$T_f = 252^{+40}_{-39} \text{ MeV} \quad (3)$$

from the first ratio, and 228^{+57}_{-228} MeV from the second. These two estimates are compatible with each other, so we can use the first number in our extraction.

This single parameter now gives predictions for the yields of the other bottomonium mesons. These are shown in Figure 1. The compatibility of the yield of the $\Upsilon(3S)$ with the 1S and 2S states is visible in this figure as the agreement between the prediction, based on $Y[\Upsilon(2S)]$, and the direct measurement. The predictions for the χ states is made under the assumption that they freeze out at the same temperature. It is generally seen in lattice computations that the P states have significantly higher widths than the S states [14]. If so, then they would freeze out later. The relative yields shown in Figure 1 are therefore upper limits of expectations.

It is also possible to examine the centrality dependence of the yields, although the errors are bound to be larger. First estimates are presented in Figure 2. We have used the double ratio reported in [9] and the value of $N_{pp}[\Upsilon(2S)]/N_{pp}[\Upsilon(1S)]$ reported there to get $Y[\Upsilon(2S)]$ as a function of centrality. The errors in the double ratio include the errors in this normalization, so adding them would result in double counting. In fact, these errors should be removed from the double ratio when we reconstruct $Y[\Upsilon(2S)]$ in different centrality bins. Unfortunately this is not possible without access to the data sample, so we just add in quadrature the statistical and systematic uncertainties in the double ratio. The near constancy of T_f could be a strong argument for thermalization, in the sense that the initial state information is forgotten.

However, it would be premature to come to this conclusion yet, in view of the large errors. Here is one of the alternative scenarios which must be examined with improved data. In peripheral collisions the separation between the comovers and fireball may be much reduced, resulting in other sources of suppression. This may result in the effective value of T_f increasing in peripheral collisions. The values shown in the figure may indicate such an effect, although it is hard to make any statistically significant inference because of the large error bars. Nevertheless, the fact that the four most peripheral bins lie on one side of the mean, and the central bins lie largely on the other side could have some significance. We believe that a re-analysis of the data along these lines by the experimental collaboration should be able to give a clearer picture.

We turn now to the preliminary data on J/ψ from the LHC. $R_{\text{PbPb}}[J/\psi]$ and its systematics have been reported by the ALICE [24], ATLAS [25] and CMS [26] collaborations. However, the comparison of the $\psi(2S)$ and J/ψ is yet to

mature. CMS [10] uses the most central 0–20% of the events to extract values

$$Y[\psi(2S)] = \frac{N_{\text{PbPb}}[\psi(2S)]}{N_{\text{PbPb}}[J/\psi]} = \begin{cases} 0.024 \pm 0.008 & (|y| \leq 1.6), \\ 0.105 \pm 0.02 & (1.5 \leq |y| \leq 2.4). \end{cases} \quad (4)$$

Both these values are boosted above the pp ratio, implying that the fireball is a little richer in $\psi(2S)$ than an equivalent system without re-interactions. The central rapidity bin contains p_T between 6.5 and 30 GeV, whereas the peripheral bin has p_T between 3 and 30 GeV. These give $T_f = 149 \pm 14$ MeV at central rapidity but 255 ± 25 MeV at larger rapidity. Since both rapidity bins are well-separated from the spectator rapidity, one does not expect this difference to be due to comover suppression effects. It is more likely that the differences arise from the different p_T acceptances. The higher p_T particles escape more easily from the fireball and therefore may not thermalize perfectly. If so, then the lower value of T_f for higher p_T cutoff simply tells us that the charmonia which are sampled are not in complete thermal equilibrium. This is borne out by the results reported by CMS on the N_{part} dependence of the double ratio for $|y| < 1.6$. Our analysis seems to show that T_f increases from about 185 MeV to 240 MeV as N_{part} increases from 35 to 310. We believe that the correct interpretation of this result is that the large p_T charmonium system goes from being less to more nearly thermalized in more central collisions, i.e., as the fireball size increases. This picture can also account for the fact that the systematics of the double ratio $R_{\text{PbPb}}[\psi(2S)]/R_{\text{PbPb}}[J/\psi]$ reported by ALICE [11] is quite the opposite, and qualitatively similar to the CMS observations for bottomonium [9]. ALICE has the advantage that it can detect J/ψ and $\psi(2S)$ with no lower cutoff on p_T . Once ALICE gives the pp normalization which allows converting the double ratio to $Y[\psi(2S)]$, a genuine T_f in the charmonium system can be extracted.

We have argued that sequential suppression of the Υ family of mesons observed in the CMS experiment at LHC [9] can be interpreted as the system being in thermal equilibrium in the fireball before freezing out. R_{PbPb} provides a simple diagnostic that the physics of PbPb collisions is different from that in pp collisions, but thermal behaviour is easier to analyze using the relative yields, Y , defined in eq. (2). The yields of $\Upsilon(nS)$ allow us to make a first estimate of a freezeout temperature, $T_f = 252 \pm 40$ MeV, which is different from the freezeout temperatures seen in hadrons made of light quarks. This is more or less constant as a function of the centrality, with perhaps a very mild tendency to drop in the most central collisions. This could mean a slightly earlier freezeout in peripheral collisions, or a mild contamination by cold nuclear matter effects. This picture predicts the yields of other bound states in the bottomonium sector. In the charmonium sector the data is still preliminary. The absence of complete data with low- p_T charmonia makes it impossible to extract its T_f now, although the situation may improve soon. We have pointed out various possible lines of investigation which probe such a simple analysis more deeply.

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